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PRODUCTION OF KALUZA-KLEIN STATES AT FUTURE COLLIDERS

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ABSTRACT

Perturbative breaking of supersymmetry in four-dimensional string theories predict in general the existence of new large dimensions at the TeV scale. Such large dimensions lie in a domain of energies accessible to particle accelerators. Their main signature is the production of Kaluza-Klein excitations which can be detected at future colliders. We study this possibility for hadron colliders (TEVATRON, LHC) and e^+e^- colliders (LEP-200, NLC-500).

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New degrees of freedom are required in any attempt of unification of the electroweak and strong interactions with gravity. Among these attempts, only superstring theory is known to provide a consistent quantum theory of gravity [1]. Strings predict two kinds of new degrees of freedom: (i) Superheavy oscillation modes whose characteristic mass scale is given by the inverse of the string tension $\alpha'^{-1/2} \sim 10^{18}$ GeV. These states are important at very short distances of order of the Planck length and modify the ultraviolet behavior of gravitational interactions. (ii) States associated to the internal compactified space whose presence is required from the fact that superstring theory in flat space is anomaly free only in ten dimensions. Usually, the size of this internal space is also made too small to give any observable effect in particle accelerators.

In this work we are interested in the possibility of having one or two large internal dimensions at a scale accessible to future experiments. The presence of such large dimensions is motivated by superstring theory with perturbative breaking of supersymmetry [2-3]. In this context, their size is inversely proportional to the scale of supersymmetry breaking which must be of order of the electroweak scale in order to protect the gauge hierarchy. In contrast to field theoretical expectations, string theory allows the existence of such large dimension(s) consistently with perturbative unification of low-energy couplings in a class of models based on orbifold compactifications [3].

Properties and implications of string models with perturbative breaking of supersymmetry, were studied in the case of minimal embedding of the standard model [4]. It was pointed out that the main signature of the large extra dimension(s) in these constructions is the appearance of a tower of excitations for the gauge bosons and higgses with the same gauge quantum numbers. These Kaluza-Klein (KK) states, characteristic of all theories with compactified dimensions, have masses

$$m_n^2 = m_0^2 + \frac{\vec{n}^2}{R^2} \ , \tag{1}$$

where R denotes the common radius of the D large internal dimensions ($D = 1$ or 2), while $\vec{n}$ is a D dimensional vector with integer entries. m_0 stands for R-independent contributions coming from the electroweak symmetry breaking. In the following, m_0 will be neglected since we are interested in values of $R^{-1} \gtrsim 200$ GeV. Another characteristic of these models is that the massive KK-modes are organized in multiplets of $N = 4$ supersymmetry, which contain one vector boson, four two-component fermions and three complex scalars. Since KK-excitations exist only for gauge bosons and higgses, they form, in the minimal case, the adjoint representation

of $SU(3)\times SU(3)_c$ [5]. In fact, the adjoint representation of the first $SU(3)$ contains in addition to the adjoint of $SU(2)\times U(1)$ two doublets with the gauge quantum numbers of the two higgses of the minimal supersymmetric extension of the standard model (MSSM). In contrast, quarks and leptons have no excitations.

All massive KK-states are unstable. They disintegrate into quarks and leptons within a short lifetime, of the order of 10^{-26} seconds, when the size of the compact dimension(s) is 1 TeV^{-1} [5].

Present experimental limits have been obtained from an analysis of the effective four-fermion operators which arise from the exchange of the massive KK-modes. Model independent bounds can be derived from the modifications of known cross-sections generated by such effective interactions [6]. A computation of these operators in orbifold models has shown that the present limits are $R^{-1}\gtrsim 185$ GeV for one large extra dimension, while $R^{-1}\gtrsim 1.4\ \mathrm{TeV}, 1.1\ \mathrm{TeV}, 1\ \mathrm{TeV}$ for two large dimensions in the case of $\mathbf{Z}_3$, $\mathbf{Z}_4$ and $\mathbf{Z}_6$ orbifolds, respectively [5]. These values leave open the exciting possibility of detecting such extra dimensions at TEVATRON, or at future colliders (LHC, LEP-200, NLC-500). Below, we compute the cross sections and asymmetries when KK-excitations are produced in these colliders.

THE COUPLINGS OF KALUZA-KLEIN STATES

For the purpose of this work, we need the values for the couplings of KK-modes to quarks and leptons. It is easy to show that, in the minimal case, these couplings are non-vanishing only for vector excitations with the quantum numbers of the standard model gauge bosons, and for scalar excitations with the quantum numbers of the two Higgs doublets [5]. The former couple through gauge interactions, while the latter through Yukawa interactions.

In the large radius limit, the tree level couplings of the massive KK-modes to matter fermions are all equal to the corresponding coupling of the lowest excitation, up to negligible α'/R^2 corrections. However, the computation of any low energy process requires the knowledge of the effective renormalized couplings. In string theory, which is ultraviolet finite, the low energy running of coupling constants can be obtained from the logarithmic infrared divergences of the corresponding on-shell scattering amplitudes [7]. This procedure is appropriate for couplings of massless particles and it becomes less clear when external legs are massive KK-modes.

One way to define the one loop renormalized couplings of the vector KK-

excitations to quarks and leptons is to extrapolate the results of ref.[5] for the coefficients of the corresponding four-fermion effective operators. These coefficients were computed in the kinematic region of external momenta smaller than the compactification scale, by summing over all one loop infrared divergent bubbles. In fact at the tree level, the four-fermion amplitude is proportional to:

$$\mathcal{A}(p,R) = \frac{g^2}{p^2} + g^2 \delta S \ , \tag{2}$$

where g is the tree level string coupling. The first term in the r.h.s. of (2) denotes the contribution of the massless gauge boson exchange (considering only the s channel), while the second term describes a contact interaction coming from the exchange of the massive KK-modes [5]. The renormalized one loop amplitude becomes [5] [1]:

$$\mathcal{A}_{ren}(p,R) = \frac{g^2(p)}{p^2} + \frac{g^4(p)}{g_U^4(p)} g^2 \delta S \ , \tag{3}$$

where $g(p)$ is the usual running coupling constant associated to the massless gauge boson, while $g_U(p)$ is the effective coupling which takes into account only the contribution to the β-function of the massless states which have KK-excitations, *i.e.* gauge boson and Higgs supermultiplets. Equation (3) suggests that the effective coupling of the massive KK-modes can be identified as:

$$\alpha^*(p) \equiv \frac{g^{*2}(p)}{4\pi} = \frac{g^4(p)}{g_U^4(p)} \alpha \ . \tag{4}$$

where $\alpha \equiv \frac{g^2}{4\pi}$. In fact, though the expression (3) has been obtained in the low momentum region ($pR \ll 1$), the coefficients of the four-fermion operators ($\sim g^{*2}(p)/m_n^2(p)$) in the effective low-energy theory should not run for scales $p^2R^2 \lesssim 1$ [2]. For scales $p^2R^2 > 1$ the coupling $g^*(p)$ should run with the scale as in an ordinary field theory. However to compute the running of $g^*(p)$ between $1/R$ and M_{GUT} we should properly take into account the (unknown) threshold

[1] In the simplest case where all matter fermions come from the same fixed point of the orbifold.

[2] A direct calculation of the electromagnetic corrections to muon beta decay shows that this is the case in the standard model of electroweak interactions [8]. More precisely, the Fermi constant $G_F \sim g^2(p)/M_W^2(p)$ does not run (it is only finitely renormalized).

effects due to the exceedingly large number of KK-excitations with masses m_n. Fortunately, we only need in our calculation the couplings $g^*(p)$ at scales $p^2 \leq m_n^2$, where the value of $g^*(p)$ is provided by the couplings $g(p)$ and $g_U(p)$, which have a well known running since they are renormalized only by the states belonging to the usual massless sectors of the theory.

Using the values for the standard model coupling constants $g_2(M_Z) = 0.656$ and $g'(M_Z) = 0.357$ and the particle content of the MSSM (*i.e.* gauge coupling unification at $M_{GUT} \sim 2 \times 10^{16}$ GeV, with $\alpha_{GUT} \sim 1/25$) one finds that in the region of energies of a few hundred GeV to a few TeV, which will be explored in future colliders, the effective coupling of the massive $SU(2)$ ($\vec{W}_n^*$) excitations $g_2^*(p)$ is of the order of 10^{-2} and can be neglected. On the other hand, the value of the effective coupling of the massive $U(1)_Y$ (B_n^*) excitations $g'^*(p)$ varies in the range 0.25-0.27 .

We should stress here that the quoted values of $g_2^*(p)$ and $g'^*(p)$ are sensitive to the unification scale. In fact, in models where M_{GUT} is shifted to $\sim 10^{18}$ GeV, $g_2^*(p)$ is no longer negligible as compared to $g'^*(p)$. In those cases, γ_n^* and Z_n^* have sizeable couplings to fermions and should both be considered in the production cross sections and asymmetries computed below, therefore modifying correspondingly the bounds on R^{-1}.

PRODUCTION AT HADRON COLLIDERS

Among the various KK-excitations of different spins, the easiest to detect at future colliders are the vectors with the quantum numbers of the electroweak $SU(2) \times U(1)_Y$ gauge bosons. Here, we shall follow the method of ref.[9], where a similar analysis was performed for the production, at hadron colliders, of Z' vector bosons present in E_6 superstring-inspired models. It was pointed out that the most efficient way of observing new gauge bosons in proton-(anti)proton collisions is to identify charged leptons $l^\pm$ in the final state. Thus, we need to compute the cross-sections for the Drell-Yan processes $pp \to l^+l^-X$ at LHC, or $p\bar{p} \to l^+l^-X$ at TEVATRON, with $l = e, \mu, \tau$. These processes receive also a contribution from the exchange of scalars with quantum numbers of the higgses, which is however suppressed by the small Yukawa couplings. Scalars, as well as the remaining vector bosons, can be produced through other processes and they are in general harder to detect.

As pointed out above, because of the accidental suppression of the effective coupling of the massive $SU(2)$ vector excitations, only the $U(1)_Y$ excitations B_n^* ($n \neq 0$) couple to leptons and contribute to Drell-Yan processes. These states have masses given in eq.(1) and they couple to matter fermions through the effective interaction:

$$g'^*(p)\bar{\psi}^k \gamma_\mu (v_k + a_k \gamma_5) \psi^k B_n^{*\mu} \ , \tag{5}$$

where k labels the different species of fermions, and $g'^*(p)$ is given in eq.(4). The values of the vector couplings v_k and the axial couplings a_k are given in table 1.

The interactions (5) lead to rates of B_n^* decays into fermions:

$$\Gamma(B_n^* \to f\bar{f}) = g'^{*2}(m_n) \ \frac{m_n}{12\pi} C_f (v_f^2 + a_f^2) \ , \tag{6}$$

while the corresponding interactions with their scalar superpartners $\tilde{f}_{R,L}$, lead to the decay rates:

$$\Gamma(B_n^* \to \tilde{f}_{\substack{R\\L}} \tilde{\bar{f}}_{\substack{R\\L}}) = g'^{*2}(m_n) \ \frac{m_n}{48\pi} C_f (v_f \pm a_f)^2 \ , \tag{7}$$

where $C_f = 1$ or 3 for color singlets or triplets, respectively. In the above computation we neglected the masses of the fermions and of their superpartners. The latter could be important for the lowest excitations, but they are model dependent. In this approximation, where both the supersymmetry and electroweak gauge symmetry breaking are neglected, the width Γ_n of the B_n^* resonances is obtained by summing only over the partial widths (6) and (7). In fact, the decay rates of B_n^* to a pair of states with gauge quantum numbers of the higgses vanish in this limit, due to the conservation of the momentum associated to the compactified dimension(s). The total width is then:

$$\Gamma_n = \frac{5}{8\pi} g'^{*2} m_n \ . \tag{8}$$

Note that if one takes into account only the decay rates into fermions, the width (8) is decreased by a factor 2/3.

The proton-(anti)proton system has a center of mass energy $\sqrt{s} = 16$ TeV and 1.8 TeV for LHC and TEVATRON, respectively. In the corresponding Drell-Yan process, the lepton pairs are produced via the subprocess $q\bar{q} \to l^+ l^- X$ of center of mass energy M. The two colliding partons take a fraction

$$x_a = \frac{M}{\sqrt{s}} \ e^y \quad \text{and} \quad x_b = \frac{M}{\sqrt{s}} \ e^{-y} \tag{9}$$

of the momentum of the initial proton (a) and (anti)proton (b), with a probability described by the quark or antiquark distribution functions $f^{(a)}_{q,\bar{q}}(x_a, M^2)$ and $f^{(b)}_{q,\bar{q}}(x_b, M^2)$. In our numerical computations, we will use the parton distribution functions given by Martin, Roberts and Stirling (set S_o) [10].

The total cross section, due to the production of the KK-excitations B^*_n, is given by:

$$\sigma = \sum_{q=\text{quarks}} \int_0^{\sqrt{s}} dM \int_{\ln(M/\sqrt{s})}^{\ln(\sqrt{s}/M)} dy\ g_q(y, M) S_q(y, M)\ , \tag{10}$$

where

$$g_q(y, M) = \frac{M}{18\pi} x_a x_b\ [f_q^{(a)}(x_a, M^2) f_{\bar{q}}^{(b)}(x_b, M^2) + f_{\bar{q}}^{(a)}(x_a, M^2) f_q^{(b)}(x_b, M^2)]\ , \tag{11}$$

and

$$S_q(y, M) = g'^{*4} \frac{1}{N} \sum_{|\vec{n}|<R\sqrt{s}} \frac{(v_q^2 + a_q^2)(v_l^2 + a_l^2)}{(M^2 - m_n^2)^2 + \Gamma_n^2 m_n^2}\ , \tag{12}$$

where the factor $1/N$ comes from the $\mathbf{Z}_N$ orbifold projection which identifies the state $|n>$ with $|\theta n>$, where $\theta = e^{\frac{2i\pi}{N}}$. In the case of two dimensions, we use the complex notation $n = n_1 + \theta n_2$ to represent the vector $\vec{n} = (n_1, n_2)$, and $\vec{n}^2 = |n|^2$. For the numerical results, we restrict our analysis to the cases of one-dimensional $\mathbf{Z}_2$ and two-dimensional $\mathbf{Z}_4$ orbifolds. In eq.(12), we have dropped the interference terms which are negligible in our case, since the KK-resonances do not overlap, as their widths (8) are small compared to their relative distance due to the coupling constant suppression. This property also allows to replace in eq.(10) the integration over M by the approximate expression $\frac{\pi}{2}\Gamma_n \frac{d\sigma}{dM}|_{M=m_n}$ for each resonance.

The cross-section (10) has to be corrected by incorporating the radiative corrections in the initial states. These corrections are usually described by a multiplicative factor K [11]. We found $K \sim 1.3$ for TEVATRON and $K \sim 1.1$ for LHC. After multiplication by K, the total cross-sections for the production of l^+l^- pairs at TEVATRON and LHC are plotted as a function of R^{-1} in figs. 1 and 2, respectively. In each figure, we present two curves for the cases of one (solid lines) and two (dotted lines) large compactified dimensions. The corresponding numbers of events are obtained after multiplication by the integrated luminosity.

At TEVATRON, the CDF collaboration has collected an integrated luminosity $\int \mathcal{L} dt = 21.4\ pb^{-1}$ during the 1992-93 running period. From the non-existence of candidate events at e^+e^- invariant mass above 350 GeV, and using the fact that the detection efficiency is rather flat and ~ 40 % in the region between 350 GeV

and 700 GeV [3], they obtain the bound [12]

$$\sigma(B_n) \cdot B(B_n \to e^+e^-) < 0.35 \; pb \; (95 \; \% \; C.L.) \tag{13}$$

From eq.(13) and the result in fig. 1 one deduces the bound $R^{-1} \lesssim 510$ GeV. We can scale the bound (13) for the mass region below 700 GeV, to the increased value of luminosity accessible in the future, $\int \mathcal{L}dt \sim 100 \; pb^{-1}$

$$\sigma(B_n) \cdot B(B_n \to e^+e^-) \lesssim 0.075 \; pb \; (95 \; \% \; C.L.) \tag{14}$$

which would translate in the (future) bound $R^{-1} \lesssim 650$ GeV. In particular, this result implies that (in view of the limit of 1.1 TeV from the analysis of effective four- fermion operators) the KK-excitations in the case of $\mathbf{Z}_4$ orbifold should not be detected at any future TEVATRON run. However, since those limits are milder for $\mathbf{Z}_2$ orbifolds, the corresponding KK-excitations could be produced if they are lighter than ~700 GeV. For LHC (fig.2), using a luminosity of $\int \mathcal{L}dt \sim 10^5 \; pb^{-1}$, and an efficiency of ~ 15 % (which would amount at 95 % C.L. to a discovery limit of 20 events [13]), one would obtain a bound

$$\sigma(B_n) \cdot B(B_n \to e^+e^-) \lesssim 2 \times 10^{-4} \; pb \; (95 \; \% \; C.L.) \tag{15}$$

i.e., $R^{-1} \lesssim 4.5$ TeV.

Above, we gave the results based on total cross sections. In the presence of enough statistics one could in principle identify different resonances, regularly spaced, by plotting the number of events as a function of the lepton pair invariant mass M. This would be a clear signal for the existence of new dimension(s). Of course, an identification of precise couplings is also required but this is much harder to achieve at hadron colliders.

PRODUCTION AT e^+e^- COLLIDERS

In contrast to the case of hadrons, in e^+e^- colliders the invariant mass of the produced fermion pairs is fixed (to a first approximation) and the presence of B_n^* resonances cannot be directly observed, unless the machine energy happens to be very close to the mass of one of the excitations, or else a scanning of energies is made.

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[3] We thank K. Maeshima for private communication on this point.

Moreover, since the presently planned energies of next linear colliders (NLC-500) are small compared to LHC, direct production of heavy vector bosons is more limited. However, e^+e^- experiments have a clean environment which allows to perform high precision measurements. These could reveal the existence of extra gauge bosons, through their virtual effects. The strategy is then to compare accurate measured quantities, such as the total cross section σ_T, the forward-backward asymmetry A_{FB}, and the ratio of hadron to lepton production, to the values predicted by the standard model. A possible disagreement could be interpreted as a signal for new physics at a higher scale. Below, we compute the deviations in σ_T and A_{FB} due to the exchange of KK-modes, at LEP-200 and NLC-500.

The total cross section for the annihilation of unpolarized electron-positron pairs e^+e^-, with a center of mass energy $\sqrt{s}$, to lepton pairs l^+l^-, through the exchange of vector bosons in the s-channel, is given by:

$$\sigma_T^0(s) = \frac{s}{12\pi} \sum_{\alpha,\beta=\gamma,Z,B_n^*} g_\alpha^2(\sqrt{s}) g_\beta^2(\sqrt{s}) \frac{(v_e^\alpha v_e^\beta + a_e^\alpha a_e^\beta)(v_l^\alpha v_l^\beta + a_l^\alpha a_l^\beta)}{(s - m_\alpha^2 + i\Gamma_\alpha m_\alpha)(s - m_\beta^2 - i\Gamma_\beta m_\beta)} \ , \tag{16}$$

where the labels α, β stand for the different neutral vector bosons γ, Z, and B_n^* with coupling constants $g_\alpha = e$, $e/(\sin\theta_w \cos\theta_w)$, and g'^*, respectively; θ_w is the weak mixing angle. With the same notation, the cross section for the forward-backward asymmetry is:

$$\sigma_{FB}^0(s) = \frac{s}{16\pi} \sum_{\alpha,\beta=\gamma,Z,B_n^*} g_\alpha^2(\sqrt{s}) g_\beta^2(\sqrt{s}) \frac{(v_e^\alpha a_e^\beta + a_e^\alpha v_e^\beta)(v_l^\alpha a_l^\beta + a_l^\alpha v_l^\beta)}{(s - m_\alpha^2 + i\Gamma_\alpha m_\alpha)(s - m_\beta^2 - i\Gamma_\beta m_\beta)} \ . \tag{17}$$

As we expect in experiments to be held at energies far from the resonances peaks, the contribution of B_n^* exchanges to the cross sections (16) and (17) will be dominated by the γ B_n^* interference terms, due to the mass suppression from the propagators. These effects are small, and they must be computed with high accuracy. It is then necessary to include radiative corrections, and in particular the bremsstrahlung effects on the initial electron and positron [14]. These are described by the convolution of (16) and (17) with radiator-functions which describe the probability of having a fractional energy loss, x, due to the initial state radiation:

$$\sigma_T(s) = \int_0^\Delta dx \sigma_T^0(s') r_T(x) \ , \tag{18}$$

and

$$A_{FB}(s) = \frac{1}{\sigma_T} \int_0^\Delta dx \sigma_{FB}^0(s') r_{FB} T(x) \ , \tag{19}$$

with

$$s' = s(1-x) \ .$$

In the above equations, Δ represents an experimental cut for the energy of emitted soft photons in bremsstrahlung processes. The radiator functions are given by [14]:

$$\begin{aligned} r_T(x) &= (1+X)yx^{y-1} + H_T(x) \\ r_{FB}(x) &= (1+X)yx^{y-1} + H_{FB}(x) \ , \end{aligned} \tag{20}$$

with:

$$\begin{aligned} X &= \frac{e^2(\sqrt{s})}{4\pi^2}[\frac{\pi^2}{3} - \frac{1}{2} + \frac{3}{2}(\log\frac{s}{m_e^2} - 1)] \\ y &= \frac{2e^2(\sqrt{s})}{4\pi^2}(\log\frac{s}{m_e^2} - 1) \\ H_T &= \frac{e^2(\sqrt{s})}{4\pi^2}[\frac{1+(1-x)^2}{x}(\log\frac{s}{m_e^2} - 1)] - \frac{y}{x} \\ H_{FB} &= \frac{e^2(\sqrt{s})}{4\pi^2}[\frac{1+(1-x)^2}{x}\frac{1-x}{(1-\frac{x}{2})^2}(\log\frac{s}{m_e^2} - 1 - \log\frac{1-x}{(1-\frac{x}{2})^2})] - \frac{y}{x} \ , \end{aligned} \tag{21}$$

where m_e is the electron mass.

We use $\sqrt{s} = 190$ GeV for LEP-200 and 500 GeV for NLC-500, together with the numerical values for the experimental cuts [14]:

$$\Delta(\mathrm{LEP}-200) = 0.770 \quad \text{and} \quad \Delta(\mathrm{NLC}-500) = 0.967 \ , \tag{22}$$

coming from the condition of removing the Z boson tail, which amounts to imposing the cut $s' \geq M_Z^2$.

In fig.3 we plot, for LEP-200, the ratios

$$R_T = \left| \frac{\sigma_T(s) - \sigma_T^{\mathrm{SM}}(s)}{\sigma_T^{\mathrm{SM}}(s)} \right| \tag{23}$$

and

$$R_{FB} = \left| \frac{A_{FB}(s) - A_{FB}^{\mathrm{SM}}(s)}{A_{FB}^{\mathrm{SM}}(s)} \right| , \tag{24}$$

where $\sigma_T^{\mathrm{SM}}(s)$ and $A_{FB}^{\mathrm{SM}}(s)$ are the standard model predictions for the total cross section and forward-backward asymmetry, respectively. The cases of $\mathbf{Z}_2$ and $\mathbf{Z}_4$ orbifolds are indicated with solid and dotted lines, respectively. The same plot for NLC-500 is exhibited on fig.4.

We can see from figs. 3 and 4 that the measurement of either of these quantities, $\sigma_T(s)$ and $A_{FB}(s)$ with a certain degree of accuracy will translate into a lower bound on R^{-1}. For instance, measuring both of them at LEP-200 with an error 1% would translate into a bound $R^{-1} \gtrsim 1.6$ TeV for the $\mathbf{Z}_2$ orbifold, and $R^{-1} \gtrsim 3.0$ TeV for the $\mathbf{Z}_4$ orbifold. At NLC-500 the bounds would become $R^{-1} \gtrsim 4.9$ TeV for the $\mathbf{Z}_2$ orbifold, and $R^{-1} \gtrsim 9.2$ TeV for the $\mathbf{Z}_4$ orbifold.

One of the reasons for building future colliders is the search for sparticles which should be seen if supersymmetry is broken at low energy. This work shows that the planned energies will also allow to partially investigate the origin of the breaking. In particular, the cross sections derived above make it clear that LHC and NLC are able to test the possibility for perturbative breaking using large extra dimension(s). Moreover, the possible detection of KK-states is very exciting since such dimensions can be implemented consistently only in superstring theory and, thus, provide a window to study a part of the massive string spectrum at low energies.

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Coupling	photon	Z	B_n
v_u	2/3	$\frac{1}{4} - \frac{2}{3}s_w^2$	$-\frac{5}{12}$
a_u	0	$-\frac{1}{4}$	$-\frac{1}{4}$
v_d	-1/3	$-\frac{1}{4} + \frac{1}{3}s_w^2$	$\frac{1}{12}$
a_d	0	$\frac{1}{4}$	$\frac{1}{4}$
v_e	-1	$-\frac{1}{4} + s_w^2$	$\frac{3}{4}$
a_e	0	$\frac{1}{4}$	$\frac{1}{4}$
v_ν	0	$\frac{1}{4}$	$\frac{1}{4}$
a_ν	0	$-\frac{1}{4}$	$-\frac{1}{4}$

Table 1.

Couplings of the matter particles to the standard gauge bosons and the B_n KK-excitations (we used $s_w \equiv \sin(\theta_w)$).

Figure Captions

Fig. 1 Total cross-section $\sigma(p\bar{p} \to l^+l^-X)$ at Tevatron as a function of the compactification scale R^{-1} (solid and dotted lines correspond to the one-dimensional Z_2 and two-dimensional Z_4 orbifolds, respectively).

Fig. 2 Total cross-section $\sigma(pp \to l^+l^-X)$ at LHC as a function of the compactification scale R^{-1} (solid and dotted lines correspond to the one-dimensional Z_2 and two-dimensional Z_4 orbifolds, respectively).

Fig. 3 Expected deviations from the standard model for the total and forward - backward asymmetry cross-sections at LEP-200, as a function of the compactification scale R^{-1} (solid and dotted lines correspond to the one-dimensional Z_2 and two-dimensional Z_4 orbifolds, respectively).

Fig. 4 Expected deviations from the standard model for the total and forward - backward asymmetry cross-sections at NLC-500 as a function of the compactification scale R^{-1} (solid and dotted lines correspond to the one-dimensional Z_2 and two-dimensional Z_4 orbifolds, respectively).

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